The parity–violating pion–nucleon coupling constant from a realistic three flavor Skyrme model

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Abstract

We study the parity–violating pion–nucleon coupling G_π in the framework of a realistic three flavor Skyrme model. We find a sizeable enhancement of $G_\pi \simeq 0.8 \dots 1.3 \cdot 10^{-7}$ compared to previous calculations in two–flavor models with vector mesons. This strangeness enhancement stems from induced kaon fields of the chiral soliton and the non-monotoneous dependence on symmetry breaking of the nucleon matrix element of the flavor singlet piece of the operator associated with these induced fields. Both features are sensitive to four quark operators including an $\bar{s}s$ pair.

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1. There has been considerable interest and controversy about the parity–violating pion–nucleon coupling constant G_π[#5] over the last years, triggered on one side by new experimental results and on the other by fresh theoretical ideas. The measurement of the anapole moment in ^{133}Cs [1], which allows to get a bound on G_π, seems to contradict the bounds from the anapole moment measured in ^{205}Tl [2] and the bound from the circular polarization asymmetry measurement of ^{18}Fl [3]. To be precise, these data are analyzed in the framework of parity–violating (pv) one–boson exchange and are thus sensitive to the products of the weak and the strong (parity–conserving) couplings. The latter are, however, sufficiently well known for the present accuracies one is dealing with. A recent paper which deals with this topic is ref. [4]. On the theoretical side, the chiral perturbation theory analysis of Kaplan and Savage [5] seems to indicate a large enhancement of the weak pion–nucleon coupling due to strangeness. More precisely, the underlying four–fermion current–current Hamiltonian has a piece of the form $(\bar{q}q)(\bar{s}s)$, which contributes sizeably to G_π. Here, q (s) denotes the light (strange) quarks. In ref. [5], numerical estimates were given based on factorization and dimensional analysis. On the other hand, the two–flavor topological chiral soliton model was used to study parity–violating meson–nucleon couplings [6, 7] and interaction regions [8], including also the $N\Delta$ and $\Delta\Delta$ vertices. This model gives a successful description of many nucleon observables as reviewed in ref. [9]. In this approach, the weak pion–nucleon coupling comes out to be very small, typically $G_\pi \simeq 0.3 \cdot 10^{-7}$. Note that in the SU(2) Skyrme model without vector mesons, the weak pion–nucleon coupling vanishes due to a particular symmetry between the currents [6]. This symmetry is absent in the presence of vector mesons or realistic three flavor version (we thus always compare to the vector meson stabilized Skyrmion when we talk about SU(2)). Clearly, in the two–flavor approch one is not sensitive to operators involving strange quarks and also strange components in the nucleon wave function. It appears therefore mandatory to extend the soliton model calculations to the three flavor case. In this letter, we calculate G_π in a realistic SU(3) Skyrme model, which gives a fair description of many observables, like the mass splittings of the low-lying $\frac{1}{2}^+$ and $\frac{3}{2}^+$ baryons, magnetic moments, hyperon deacys and many others (for a review, see ref. [10]). Within this approach, we can quantify the role of the four–quark operators involving strange quark pairs as well as the role of strangeness in the nucleons' wave function.

2. The three flavor Skryme model is defined by the Lagrange density

$$
\begin{aligned}
\mathcal{L} &= -\frac{f_\pi^2}{4}\text{tr}\ (\alpha_\mu \alpha^\mu) + \frac{1}{32e^2}\text{tr}\left([\alpha_\mu,\alpha_\nu][\alpha^\mu,\alpha^\nu]\right)\\
&\quad +\text{tr}\left\{\mathcal{M}\left[\beta'\left(\alpha_\mu\alpha_\mu U + U^\dagger\alpha_\mu\alpha_\mu\right) + \delta'\left(U + U^\dagger - 2\right)\right]\right\}\ ,
\end{aligned}
\tag{1}
$$

where 'tr' denotes the trace in flavor space, $f_\pi \simeq 93\,\text{MeV}$ the pion decay constant and e is a dimensionless number (the so–called Skyrme parameter), which can be determined from a best fit to baryon observables. Typically, $e \simeq 4$. The chiral field U, which parametrizes the Goldstone boson octet, is contained in $\alpha_\mu = \partial_\mu U U^\dagger$ and $\beta_\mu = U^\dagger \partial_\mu U$. It has been demonstrated before that for realistic three flavor calculations, one has to include symmetry breaking terms of the form given in eq.(1). The parameters associated

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[#5]We reserve the symbol f_π for the weak pion decay constant. In many papers, G_π is denoted by h_π^1.

with the symmetry breaking are determined via

$$m_\pi^2 = \frac{4\delta'}{f_\pi^2}, \quad m_K^2 = \frac{2(1+x)\delta'}{f_K^2} \quad \text{and} \quad \left(\frac{f_K}{f_\pi}\right)^2 = 1 + 4\beta'\frac{1-x}{f_\pi^2}, \tag{2}$$

as $\mathcal{M} = y\lambda^3 + T + xS$ is proportional to the current quark mass matrix (for a detailed discussion, see ref. [10]), m_π (m_K) is the pion (kaon) mass and $f_K = 1.2 f_\pi$ the kaon decay constant. Here we ignore isospin breaking, *i.e.* we set $y = 0$. Finally the Wess–Zumino–Witten term has to be added to the action

$$\Gamma_{\rm WZ} = -\frac{iN_C}{240\pi^2} \int_{M_5} d^5x\, \epsilon_{\mu\nu\rho\sigma\tau}\, \alpha^\mu \alpha^\nu \alpha^\rho \alpha^\sigma \alpha^\tau\;, \tag{3}$$

with N_C the number of colors. The classical soliton is obtained from the hedgehog *ansatz*

$$U_0(\boldsymbol{r}) = \left(\begin{array}{cc|c} \exp\left(i\boldsymbol{\tau}\cdot\hat{\boldsymbol{r}}F(r)\right) & \begin{array}{c} | \\ | \end{array} & \begin{array}{c} 0 \\ 0 \end{array} \\ \hline 0 \quad 0 & | & 1 \end{array}\right). \tag{4}$$

Here, $F(r)$ is the so–called chiral angle which is obtained from minimization of the soliton mass. Since the classical soliton has neither good isospin (hypercharge) nor spin angular momentum, one has to perform an adiabatic rotation to obtain the physical baryon states. The pertinent collective coordinates are introduced via [11]

$$U(\boldsymbol{r},t) = A^\dagger(t)\, {\rm e}^{iZ(\boldsymbol{r})}\sqrt{U_0}\, {\rm e}^{2iZ(\boldsymbol{r})}\sqrt{U_0}\, {\rm e}^{iZ(\boldsymbol{r})}A(t)\;, \quad Z(\boldsymbol{r}) = \frac{1}{2}\left(\begin{array}{c|c} 0 & K(\boldsymbol{r}) \\ \hline K^\dagger(\boldsymbol{r}) & 0 \end{array}\right), \tag{5}$$

and the time dependence of the collective coordinates A is measured by the angular velocities Ω^a:

$$A^\dagger\frac{d}{dt}A = \frac{i}{2}\sum_{a=1}^{8}\lambda^a\Omega^a. \tag{6}$$

Kaon fields $K(\boldsymbol{r})$ are induced by a linear coupling to the angluar velocities. This linear coupling term is contained in the Wess–Zumino action. An appropriate *ansatz* to the corresponding variational equations is

$$K(\boldsymbol{r}) = W(r)\boldsymbol{\tau}\cdot\hat{\boldsymbol{r}}\,\Omega_K\;, \quad \Omega_K = \frac{1}{2}\left(\begin{array}{c} \Omega_4 - i\Omega_5 \\ \Omega_6 - i\Omega_7 \end{array}\right). \tag{7}$$

These induced fields have a nonvanishing overlap

$$\langle K|K_0\rangle \propto \int_0^\infty r^2 dr\, W(r)\,\mathcal{G}\,\frac{\sin\frac{F}{2}}{1-\cos\frac{F}{2}} \tag{8}$$

with the would–be zero–mode $W_0(r) = (1 - \cos\frac{F}{2})/\sin\frac{F}{2}$. Demanding that there is no double counting of W_0 for the strange moment of inertia requires the constraint $\langle K|K_0\rangle \equiv 0$ with the metric [12]

$$\mathcal{G} = \left\{2f_\pi^2 + \frac{1}{2e^2}\left(F'^2 r^2 + 2\sin^2 F\right)\right\}(1 - \cos F)\;. \tag{9}$$

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However, we will consider various forms for $\mathcal{G}$ to illuminate the effects of the induced kaon fields. One finally ends up with the Hamiltonian for the collective coordinates

$$H = M_{\mathrm{cl}} + \frac{1}{2\alpha^2}\,J(J+1) + \frac{1}{2\beta^2}\left[C_2 - J(J+1) - \frac{3}{4}\right] + \frac{1}{2}\gamma\,(1 - D_{88}) \tag{10}$$

The classical mass M_{cl}, the moment of inertia α^2 and the symmetry breaker γ are functionals of the chiral angle F only. The induced kaon field $W(r)$ is obtained by maximizing the strange moment of inertia β^2 subject to the above constraint. $C_2 = \sum_{a=1}^{8} R_a R_a$ is the quadratic Casimir operator defined in terms of the right SU(3) generators R_a. The latter are the momentum operators conjugate to the angular velocities, $i.e.$ $R_a = -\partial L/\partial\Omega_a$. Finally the adjoint representation

$$D_{ab} = \frac{1}{2}\mathrm{tr}\left(\lambda_a A\lambda_b A^\dagger\right) \tag{11}$$

has been introduced. Exact diagonalization of H yields the wave functions of the low lying $\frac{1}{2}^+$ and $\frac{3}{2}^+$ baryons in the space of the collective coordinates. Having defined the model, we proceed to calculate the pv pion–nucleon coupling constant.

3. In the standard model the parity–violating interaction is contained in the coupling of the heavy vector bosons $(W_\mu^\pm, Z_\mu)$ to the charged and neutral quark currents, denoted J_{ch}^μ and J_{n}^μ, respectively,

$$\mathcal{L}_{\mathrm{PV}} = \frac{g_2}{\sqrt{8}}\left(W_\mu^+ J_{\mathrm{ch}}^\mu + W_\mu^- (J_{\mathrm{ch}}^\mu)^\dagger\right) - \frac{g_2}{2\cos\Theta_W}Z_\mu J_{\mathrm{n}}^\mu \ . \tag{12}$$

Here, Θ_W is the weak mixing angle (Weinberg angle). These currents are identified with combinations of the vector and axial–vector currents [13] of the three flavor Skyrme model (for $N_C = 3$) (for more details, see e.g. ref. [7])

$$\begin{aligned}
V_\mu^a(A_\mu^a) &= -\frac{i}{2}f_\pi^2\,\mathrm{tr}\left\{Q^a\left(\alpha_\mu \mp \beta_\mu\right)\right\} - \frac{i}{8e^2}\mathrm{tr}\left\{Q^a\left([\alpha_\nu,[\alpha_\mu,\alpha_\nu]] \mp [\beta_\nu,[\beta_\mu,\beta_\nu]]\right)\right\} \\
&\quad -\frac{1}{16\pi^2}\epsilon^{\mu\nu\rho\sigma}\mathrm{tr}\left\{Q^a\left(\alpha_\nu\alpha_\rho\alpha_\sigma \pm \beta_\nu\beta_\rho\beta_\sigma\right)\right\} \\
&\quad -i\beta'\mathrm{tr}\left\{Q^a\left(\{U\mathcal{M} + \mathcal{M}U^\dagger, \alpha_\mu\} \mp \{\mathcal{M}U + U^\dagger\mathcal{M}, \beta_\mu\}\right)\right\},
\end{aligned} \tag{13}$$

where the $Q^a = (\frac{1}{3}, \frac{\lambda^1}{2}, \ldots, \frac{\lambda^8}{2})$ denote the Hermitian nonet generators. Elimination of the gauge bosons in the small momentum limit $(g_2^2/4m_W^2 = \sqrt{2}G_F)$ yields the current–current interaction [6]

$$\mathcal{L}_{\mathrm{PV}} = \sqrt{2}G_F\left\{\sum_{i=1}^{2} V_\mu^i A^{i\mu} - \left[\cos(2\Theta_W)\left(V_\mu^3 + \frac{1}{\sqrt{3}}V_\mu^8\right) - \frac{1}{2}V_\mu^0\right]\left[\frac{1}{2}V^{0\mu} - A^{3\mu} - \frac{1}{\sqrt{3}}A^{8\mu}\right]\right\} \ . \tag{14}$$

Note that with $J_\mu^0 \to (2/3)J_\mu^0$ and $J_\mu^8 \to J_\mu^0/\sqrt{3}$ ($J_\mu = V_\mu, A_\mu$) one recovers the two–flavor result of ref. [7]. As noted before, for the pure SU(2) Skyrmion, the symmetry between the axial, vector and isoscalar currents leads to a vanishing G_π. This accidental symmetry is broken when one includes vector mesons or works in a realistic three flavor

version as done here. Furthermore, as argued in ref. [6], the so–called strong interaction enhancement factors are contained in the non–perturbative soliton model currents.

The pv pion–nucleon coupling constant G_π is defined via

$$\mathcal{L}_{\pi N}^{\mathrm{PV}} = -\sqrt{\frac{1}{2}}\, G_\pi\, \bar{\Psi}_N\, (\boldsymbol{\tau} \times \boldsymbol{\pi})_3\, \Psi_N \ , \tag{15}$$

where Ψ_N denotes the nucleon spinor and $\boldsymbol{\tau}$ the conventional Pauli isospin matrices. From the above current–current interaction, we have to extract the terms linear in the pion field π. This is done by considering pionic fluactuations around the chiral field U (see e.g. ref. [7])

$$U \to \eta\, U\, \eta \quad , \qquad \eta = \exp\left(\frac{i}{2 f_\pi} \boldsymbol{\tau} \cdot \boldsymbol{\pi}\right) \approx 1 + \frac{i}{2 f_\pi} \boldsymbol{\tau} \cdot \boldsymbol{\pi} \tag{16}$$

yielding the Kroll–Ruderman relations

$$V_\mu^a \to V_\mu^a + \frac{1}{2 f_\pi} F^{iab} \pi_i A_\mu^b \ , \qquad A_\mu^a \to A_\mu^a + \frac{1}{2 f_\pi} F^{iab} \pi_i V_\mu^b \tag{17}$$

up to minor contributions form the kinetic symmetry breaking term (β', β'').[#6] The interesting structure constants for the problem under consideration are $F^{i8b} = F^{i0b} = 0$ and $F^{ijb} = 2\epsilon_{ijb}$. The coupling constant G_π is now obtained from the matrix element of the term linear in π in $\mathcal{L}_{\mathrm{PV}}$ between nucleon states (with spin up):

$$\begin{aligned} G_\pi \;=\; & \frac{G_F}{f_\pi} \langle p \uparrow | \int d^3 r \Big\{ 2 \sin^2 \Theta_W \left(A_\mu^3 A^{\mu +} + V_\mu^3 V^{\mu +} \right) + \frac{1}{2} V_\mu^0 V^{+\mu} \\ & - \cos(2\Theta_W) \left[A_\mu^+ \left(\frac{1}{\sqrt{3}} A^{8\mu} - \frac{1}{2} A^{0\mu} \right) + \frac{1}{\sqrt{3}} V_\mu^8 V^{+\mu} \right] \Big\} |n \uparrow\rangle. \end{aligned} \tag{18}$$

Here the superscript '+' denotes the standard spherical component $V^+ = (V_1 + iV_2)/\sqrt{2}$ and similarly $A^+ = (A_1 + iA_2)/\sqrt{2}$. The above matrix element is understood to be taken in the space of the collective coordinates A as the nucleon states are eigenstates of the collective Hamiltonian H. The relevant operators are obtained by subsituting the meson fields in the covariant expressions for the currents. To be precise ($i, j, k = 1, 2, 3$; $\alpha, \beta = 4, .., 7$)

$$\begin{aligned} V_i^a \;=\; & V_1(r) \epsilon_{ijk} x_j D_{ak} + \frac{\sqrt{3}}{2\alpha^2} B(r) \epsilon_{ijk} x_j R_k D_{a8} - \frac{1}{\beta^2} V_2(r) \epsilon_{ijk} x_j d_{k\alpha\beta} D_{a\alpha} R_\beta \\ & + V_3(r) \epsilon_{ijk} x_j D_{88} D_{ak} + V_4(r) \epsilon_{ijk} x_j d_{k\alpha\beta} D_{a\alpha} D_{8\beta} \quad (a = 1, .., 8) \end{aligned} \tag{19}$$

$$V_0^a = \frac{\sqrt{3}}{2\alpha^2} B(r) D_{a8} + \frac{1}{\alpha^2} V_7(r) D_{ai} R_i + \frac{1}{\beta^2} V_8(r) D_{a\alpha} R_\alpha \tag{20}$$

[#6]These small corrections could be taken into account but play no role for the later results.

$$V_0^0 = B(r), \qquad V_i^0 = \frac{1}{\alpha^2} B(r) \epsilon_{ijk} x_j R_k. \tag{21}$$

for the vector current and

$$
\begin{aligned}
A_i^a &= [A_1(r)\delta_{ij} + A_2(r)\hat{x}_i\hat{x}_j] D_{aj} - \frac{1}{\beta^2} [A_3(r)\delta_{ij} + A_4(r)\hat{x}_i\hat{x}_j] d_{j\alpha\beta} D_{a\alpha} R_\beta \\
&\quad + [A_5(r)\delta_{ij} + A_6(r)\hat{x}_i\hat{x}_j] D_{aj} D_{88} + [A_7(r)\delta_{ij} + A_8(r)\hat{x}_i\hat{x}_j] D_{a8} D_{8j} \\
&\quad + [A_9(r)\delta_{ij} + A_{10}(r)\hat{x}_i\hat{x}_j] d_{j\alpha\beta} D_{a\alpha} D_{8\beta}
\end{aligned}
\tag{22}
$$

for the octet axial vector current. In the latter case the time component vanishes while the axial singlet current is omitted as it only acquires a negligibly small contribution from the β' type symmetry breaker. In the above equations the angular velocities have been eliminated in favor of the SU(3) generators R_a. The radial functions $V_1, B, .., A_1, ..., A_{10}$ can be extracted from the literature (see e.g. refs. [13],[14]).

The collective matrix elements are evaluated by insertion of a complete set of states, for example

$$\langle p \uparrow | A_\mu^+ A^{8\mu} | n \uparrow \rangle = \sum_{B;I,J} \langle p \uparrow | A_\mu^+ | B; I, J \rangle \langle B; I, J | A^{8\mu} | n \uparrow \rangle \tag{23}$$

with the sum running over all baryon eigenstates B (with isospin I and spin J) of the collective Hamiltonian H to which the nucleon can couple. In the absense of flavor symmetry breaking these would be nucleon states in higher dimensional SU(3) representations. With the inclusion of the γ term in the Hamiltonian H these intermediate states are distorted as well.

4. We are now in the position to present the results. As input, we use $f_\pi = 93 \text{ MeV}$, $f_K = 113 \text{ MeV}$, $m_\pi = 138 \text{ MeV}$, $m_K = 495 \text{ MeV}$, $G_F = 1.16 \cdot 10^{-5} \text{ GeV}^{-2}$ and $\sin^2 \Theta_W = 0.23$. The Skyrme parameter is varied in the range $e = 4.0 \ldots 4.5$, as our central value we use $e = 4.25$ which reasonably reproduces the baryon spectrum. The results for a large number of hyperon properties for these parameters can be found in the literature [10]. In fig. 1 we show the weak pion–nucleon coupling constant as a function of the kaon mass, i.e. as a function of the symmetry breaking. For the physical value of m_K, we get

$$G_\pi = \{0.8, 1.3\} \cdot 10^{-7} \quad \text{for} \quad e = \{4.0, 4.5\}, \tag{24}$$

which is considerably bigger than the SU(2) generalized Skyrme result of $0.2 \ldots 0.3 \cdot 10^{-7}$ [6, 7].[#7] However, as one freezes out the kaon degrees of freedom, the value of G_π approaches zero as indicated in fig. 1. The values given in eq.(24) are in fair agreement with most recent quark model calculations in which $G_\pi = 2.0 \ldots 2.7 \cdot 10^{-7}$ [15, 16]. The typical range of values in the quark model calculations is $G_\pi = 0 \ldots 3 \cdot 10^{-7}$ [15]. This large enhancement of the weak pion–nucleon coupling compared to the SU(2) calculations is largely due to the induced kaon fields as shown in fig. 2. Given in that figure are the proton matrix elements of the operators $O_1 = d_{3\alpha\beta} D_{8\alpha} R_\beta D_{33}$ and $O_0 = D_{83} D_{33}$. Assuming isospin invariance these operators appear, for example, when evaluating the matrix element $\langle p | A_i^+ A_i^8 | n \rangle$, *cf.* eq. (22). While the coefficient of the former contains the

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[#7]Remember that in the SU(2) Skyrme model without vector mesons, one has $G_\pi = 0$ [6].

induced kaon fields, the latter one is given entirely in terms of the classical fields. A typical example for a coefficient of O_1 would be $A_1 \times A_3$ while the one of O_0 is just like $(A_1)^2$. Again, for the explicit expressions of these radial functions we refer to the literature [13]. We note that the matrix element of O_1 is not only sizeably bigger than that of O_0 but also it is almost constant under symmetry breaking while O_0 has dropped to half of its flavor symmetric value at $\omega^2 = 3\gamma\beta^2/2 = 0$ for realistic symmetry breaking $\omega^2 \approx 5$. Actually the matrix element of O_1 has a positive slope at $\omega^2 = 0$. This underlines the importance of the induced components for the observable under consideration. These results look at first glance surprising since for most observables, the induced kaon fields play only a minor role [10]. However, a closer look at table 1 of ref. [13] reveals that the isoscalar magnetic moments $(n+p, \Lambda, \Sigma^0)$ are given up to 50% from the induced components (the V_2 contribution in that table). Furthermore, the isoscalar part of V_2 contains matrix elements of the operator $d_{3\alpha\beta} D_{8\alpha} R_\beta$. These contributions are not SU(3) symmetric by themselves but in the case of the magnetic moments tend to restore the SU(3) symmetry [13]. For the weak pion–nucleon coupling constant, this means that four–fermion operators of the type $(\bar{q}q)(\bar{s}s)$ contribute significantly because they are enhanced by the induced kaon fields. More quantitatively, setting $W(r) = 0$ reduces the value of G_π from $1.1 \cdot 10^{-7}$ to $0.5 \cdot 10^{-7}$ for $e = 4.25$. This has to be contrasted to the small induced kaon contribution to the strange moment of inertia β^2. Using $e = 4.25$ we find $\beta^2 = 0.27/\text{GeV}$ setting $W(r)$ equal to zero this is only moderately reduced to $0.24/\text{GeV}$. This underlines the statement that for three flavor calculations, one has to work with realistic symmetry breaking. Formally, such effects are suppressed in the large color limit but for the real world with $N_C = 3$ they have to be taken into account as clearly demonstrated in this calculation. We stress that the large enhancement of G_π compared to the two–flavor case is not an effect of a large strangeness component in the nucleons' wave function, but rather the suppression of the four–fermion operators involving strange quark pairs in the two–flavor case. We have furthermore studied the dependence on the metric $\mathcal{G}$. For $e = 4.25$ and $m_K = 495\,\text{MeV}$, setting $\mathcal{G} = 0\,(1 - \cos F)$ leads to $G_\pi = 2.6\,(1.9) \cdot 10^{-7}$, and the induced kaon fields contribute as much as 40% (25%) to β^2. It is therefore mandatory to account for the correct metric $\mathcal{G}$ as given in eq.(9). Let us finally also note that similarly to G_π one can evaluate the widths of non–leptonic hyperon decays in the standard model starting from the current–current interaction. In a recent study [17] it has been observed that also for these observables an important enhancement is due to the inclusion of strange degrees of freedom.

5. In summary, we have presented the results of a three flavor Skyrme model calculation for the parity–violating pion nucleon coupling constant. The resulting number of $G_\pi = 0.8 \ldots 1.3 \cdot 10^{-7}$ is in fair agreement with most recent quark model calculations and is considerably bigger than previously found values in two–flavor Skyrme type models. This large enhancement underlines the importance of four–fermion operators involving strange quark pairs. For the parity–violating vector meson couplings to the nucleon, one does not expect such dramatic changes compared to the two–flavor calculation since no unnatural suppression is involved in these cases. Nevertheless, a detailed study in a realistic three flavor vector meson Skyrme model should be carried out. In addition, one can also investigate the parity–violating vertices including the Δ resonance and the extension of the pertinent interaction regions.

Acknowledgements

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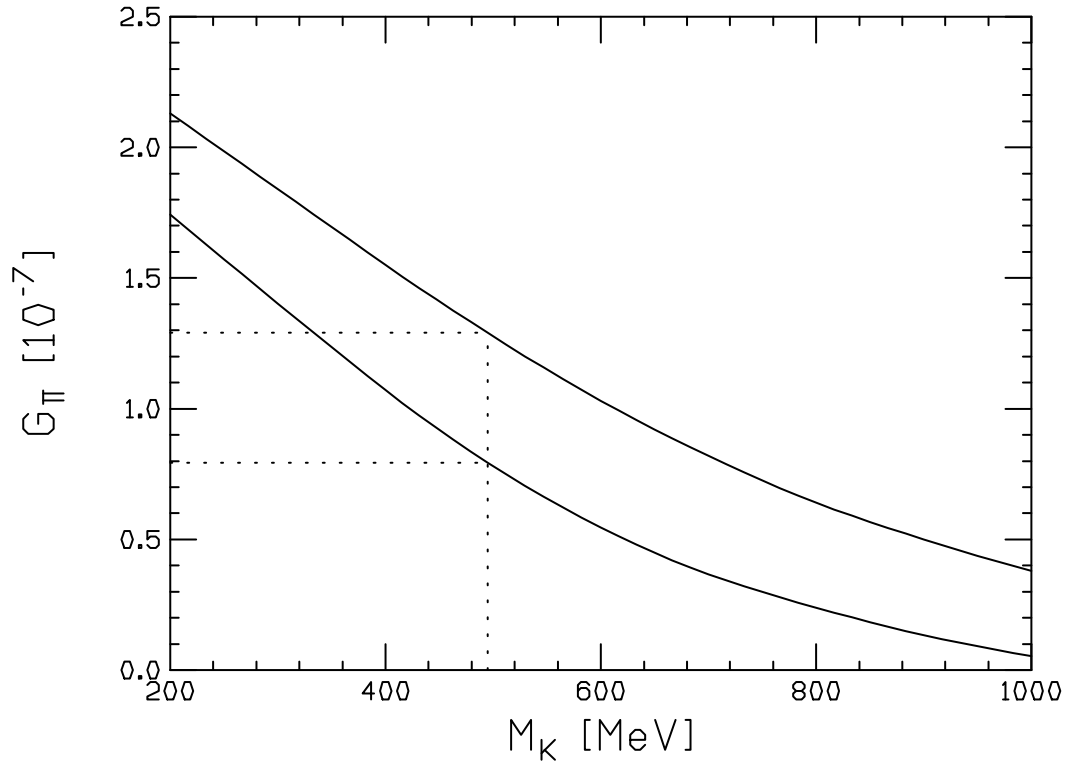
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Figures

Figure 1: The weak pion–nucleon coupling constant as a function of the kaon mass. Upper (lower) solid curve: $e = 4.5\,(4.0)$. For $m_K \to \infty$, one recovers the SU(2) result. The dotted lines refer to the physical kaon mass.
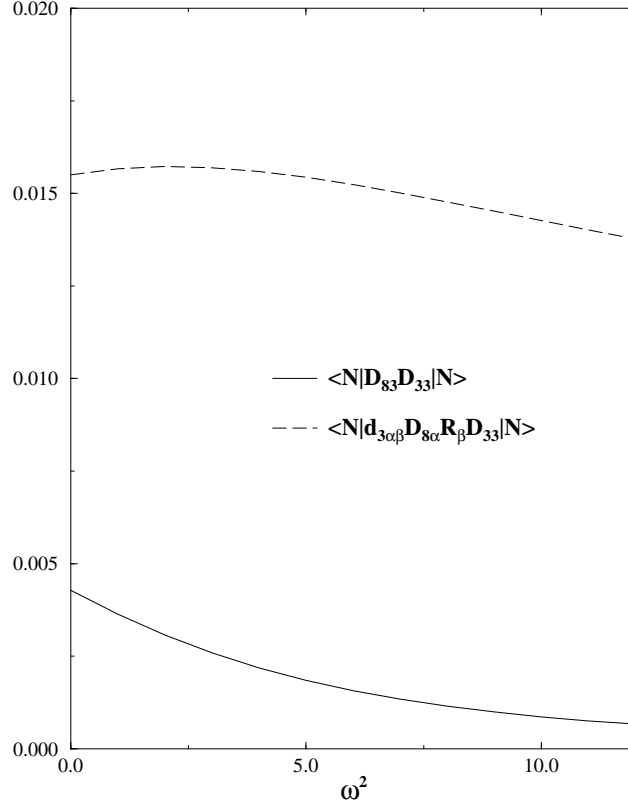
SU(3) Matrix Elements

Figure 2: SU(3) matrix elements as a function of the symmetry breaking. The solid (dashed) line refers to the classical (induced) kaon fields. For realistic symmetry breaking, $\omega^2 = \frac{3}{2}\gamma\beta^2 \simeq 5$.